Resistivity in the spin-gap state of the t-J model

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Being motivated by recent experimental data on YBaCuO, we calculate dc resistivity ρ in the spin-gap state of charge-spin-separated t-J model by using a massive gauge theory of holons and spinons. The result shows $\rho(T)$ deviates downward from the T-linear behavior below the spin-gap on-set temperature T^* as $\rho(T) \propto T\{1-c(T^*-T)^d\}$ where the mean field value of d is 1/2. To achieve smooth deviation from the T-linear behavior, one needs $d > 1$. The deviation becomes reduced with increasing hole doping.

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Many quasi-two-dimensional cuprates with high superconducting transition temperature T_c exhibit anomalous metallic behavior above T_c in Hall coefficient, magnetic susceptibility, etc. [1], which calls for a new theoretical explanation, probably in a framework beyond the conventional Fermi-liquid theory. Anderson [2] pointed out that charge-spin separation (CSS) phenomenon may be a possible explanation.

In the t-J model of strongly-correlated electrons, the CSS is naturally described by the slave-boson (SB) (or slave-fermion) mean field theory (MFT). When one incorporates fluctuations of MF's, they behave as gauge fields coupled to holons and spinons and the system possesses a U(1) gauge symmetry. The CSS can be interpreted as a deconfinement phenomenon of this gauge theory. The system is expected to have a confinement-deconfinement (CD) phase transition and the transition temperature $T_{\rm CSS}$ can be estimated [3], below which the CSS takes place.

The observed T-linear behavior of dc resistivity $\rho(T)$ [4] has been taken as one of signals of the universal anomaly in the metallic phase of high-T_c cuprates. Lee and Nagaosa [5] showed $\rho(T) \propto T$ for fermions and bosons interacting with a massless gauge field. This system has some relation to the uniform RVB MFT of the t-J model in the SB representation.

Recent experiments on YBaCuO by Ito et al. [6] and others reported that $\rho(T)$ deviates downward from the T-linear behavior below certain temperature $T^*(> T_c)$. This T^* coincides with the temperature determined by NMR and neutron experiments [7] at which a spin gap starts to develop. So it is quite interesting to calculate ρ in the spin-gap state of the SB t-J model.

The effective theory used in Ref. [5] is inadequate for this purpose, since it assumes no spin gap, containing only one gauge field associated with the hopping amplitudes of holons and spinons. In Ref. [3], we introduced multiple gauge fields, and argued that, when a spin gap develops, these gauge fields become massive due to a gauge version of Anderson-Higgs mechanism. Below we use an effective gauge field theory of holons and spinons emerged from these considerations and calculate ρ for the region $T_c < T < T^*$ to obtain $\rho \propto T\{1-c(T^*-T)^d\}$, near T^* where d is the critical exponent of gauge-boson mass, $m_A \propto (T^*-T)^d$. This reduction reflects the fact that m_A^2 suppresses gauge-field fluctuations and inelastic scatterings between gauge bosons and holons and spinons. The MFT gives $d = 1/2$, but to obtain a more reliable ρ, one needs to calculate d by renormalization-group method.

The SB t-J Hamiltonian is given by

$$H = -t\sum_{x,i,\sigma}\left(b^\dagger_{x+i}f^\dagger_{x\sigma}f_{x+i\sigma}b_x + \text{H.c.}\right) - \frac{J}{2}\sum_{xi}\left|f^\dagger_{x\uparrow}f^\dagger_{x+i\downarrow} - f^\dagger_{x\downarrow}f^\dagger_{x+i\uparrow}\right|^2. \quad (1)$$

$f_{x\sigma}$ is the fermionic spinon operator with spin $\sigma(=\uparrow,\downarrow)$ at site x of a 2d lattice [8], and b_x is the bosonic holon operator. The direction index $i(=1,2)$ is used also as unit vectors.

The MFT is obtained by decoupling both t and J terms. For definiteness we follow Ubbens and Lee [9],

$$H_{\rm MF} = \sum_{xi}\left\{\frac{3J}{8}|\chi_{xi}|^2 + \frac{2}{3J}|\lambda_{xi}|^2\right\} - \sum_{xi}\left\{\chi_{xi}\left(\frac{3J}{8}\sum_\sigma f^\dagger_{x+i\sigma}f_{x\sigma} + tb^\dagger_{x+i}b_x\right) + \text{H.c.}\right\} - \frac{1}{2}\sum_{x,i,\sigma}\left\{\lambda_{xi}\left(f^\dagger_{x\uparrow}f^\dagger_{x+i\downarrow} - f^\dagger_{x\downarrow}f^\dagger_{x+i\uparrow}\right) + \text{H.c.}\right\}. \quad (2)$$

The partition function $Z(\beta)$ $[\beta \equiv (k_BT)^{-1}]$ in path-integral formalism is given by integrating out two MF's,

the hopping amplitude, χ_{xi}, and RVB amplitude, λ_{xi} defined on the link $(x, x+i)$ [10];

$$Z = \int [db][df][d\chi][d\lambda] \exp A,$$
$$A \equiv \int_0^\beta d\tau \left\{ -\sum (b_x^\dagger \dot{b}_x + \sum f_{x\sigma}^\dagger \dot{f}_{x\sigma}) - H_{\rm MF} \right\}. \quad (3)$$

In the CSS state ($T < T_{\rm CSS}$), $\langle \chi_{xi} \rangle \neq 0$. The spin-gap state may realizes in CSS and is characterized by a condensation of λ_{xi}, $\langle \lambda_{xi} \rangle \neq 0$. Let us parameterize $\chi_{xi} = \chi U_{xi}$, $U_{xi} \equiv \exp(iA_{xi})$, $\lambda_{xi} = (-)^i \lambda V_{xi}$, $V_{xi} \equiv \exp(iB_{xi})$, assuming uniform RVB. If one ignores phase fluctuations by setting $A_{xi} = B_{xi} = 0$, spinon excitations has the energy $E(\boldsymbol{k})$, $E^2(\boldsymbol{k}) = \{(3J\chi/4)\sum_i \cos k_i - \mu_F\}^2 + \{\lambda \sum_i (-)^i \cos k_i\}^2$, where μ_F is the chemical potential to ensure $\langle f_{x\uparrow}^\dagger f_{x\uparrow} + f_{x\downarrow}^\dagger f_{x\downarrow} \rangle = 1 - \delta$ (δ is doping). We introduce also μ_B for $\langle b_x^\dagger b_x \rangle = \delta$. There appears a spin gap $\lambda(\cos k_1 - \cos k_2)$.

Under the gauge transformation, $b'_x = \exp(i\theta_x) b_x$, $f'_{x\sigma} = \exp(i\theta_x) f_{x\sigma}$, the phases of MF's transform as $A'_{xi} = A_{xi} + \theta_x - \theta_{x+i}$, $B'_{xi} = B_{xi} + \theta_x + \theta_{x+i}$. Their behavior can be studied by the effective lattice gauge theory $A_{\rm LGT}(U, V)$ that is obtained by integrating over b_x and $f_{x\sigma}$, e.g., by hopping expansion. For $T^* < T < T_{\rm CSS}$, $A_{\rm LGT}$ contains only U's interacting via the conventional gauge couplings like $\chi^4 U_{x\,2} U_{x+2\,1} U^\dagger_{x+1\,2} U^\dagger_{x\,1}$. Their quadratic terms for A_{xi} show that A_{xi} behaves as a massless gauge field. For $T_c < T < T^*$, λ develops and new couplings like $\chi^2 \lambda^2 V^\dagger_{x+1\,2} U_{x+2\,1} V_{x\,2} U_{x\,1}$ are generated. This gives rise to a mass term of A_{xi}, $m_A^2 A_{xi}^2$ with $m_A^2 = \chi^2 \lambda^2$. Also, the quadratic B_{xi} are also massive.

The holon part $A^B_{\rm eff}$ of the effective low-energy continuum field theory $A_{\rm eff}$ is obtained as

$$A^B_{\rm eff} = \int d\tau d^2x \left\{ -\bar{b}\,(\partial_\tau - \mu_B)\, b - \frac{1}{2m_B} |D_i b|^2 \right\}, \quad (4)$$

where $D_i \equiv \partial_i - igA_i$ (we introduced the gauge coupling constant $g = 1$ for convenience), and $(2m_B)^{-1} = \chi\ t\ a^2$ (a is the lattice spacing). The spinon part $A^F_{\rm eff}$ may be written as

$$A^F_{\rm eff} \simeq \int d\tau d^2x \left\{ -\bar{f}\,(\partial_\tau - \mu_F)\, f - \frac{1}{2m_F} |D_i f|^2 \right\}$$
$$- \int d\tau d^2k \left\{ \Delta_{\rm SG}(\boldsymbol{k}) \bar{f}_\uparrow(\boldsymbol{k}, \tau) \bar{f}_\downarrow(-\boldsymbol{k}, \tau) + {\rm H.c.} \right\}, \quad (5)$$

with $(2m_F)^{-1} = 3J\chi a^2/8$. Here we introduced the continuum version of the spin gap

$$\Delta_{\rm SG}(\boldsymbol{k}) \equiv \pi(1-\delta)\hat{\lambda}(T) \frac{k_x^2 - k_y^2}{k_F^2}, \quad (6)$$

where $k_F (\equiv \sqrt{2m_F\mu_F})$ is the Fermi momentum of spinons and $k_F \simeq \sqrt{2\pi(1-\delta)}/a$ at $T \ll T_F$ (T_F is the Fermi temperature). In (6), we use the renormalized spin gap $\hat{\lambda}(T)$ defined as $\langle \lambda_{xi} \rangle = (-)^i \hat{\lambda}(T)$, instead of its MF value λ to take into account the effect of phase fluctuations of λ_{xi} effectively. Below we calculate the resistivity of the system $A_{\rm eff} = A^B_{\rm eff} + A^F_{\rm eff}$ [11].

The propagator of the gauge field, $D_{ij}(\boldsymbol{x}, \tau) \equiv \langle A_i(\boldsymbol{x}, \tau) A_j(0, 0) \rangle$, is generated by fluctuations of spinons and holons, i.e., $D_{ij} = (\Pi_F + \Pi_B)^{-1}_{ij}$, where

$$\Pi_{F,B\ ij}(\boldsymbol{x}, \tau), \equiv -\langle J_{F,B\ i}(\boldsymbol{x}, \tau) J_{F,B\ j}(0, 0) \rangle_{\rm 1PI}$$
$$+ \delta_{ij} \delta(\boldsymbol{x}) \delta(\tau) n_{F,B}, \quad (7)$$

representing one-particle-irreducible (1PI) diagrams of spinon and holon loops. $J_{Fi} \equiv (2m_F)^{-1} \sum_\sigma \{ i\bar{f}_\sigma \partial_i f_\sigma + {\rm H.c.}\}$ and $J_{Bi} \equiv (2m_B)^{-1} \{ i\bar{b}\partial_i b + {\rm H.c.}\}$ are currents coupled to A_i, and $n_F = (1-\delta)/a^2$ and $n_B = \delta/a^2$. In the Coulomb gauge, the propagator at momentum $\boldsymbol{q}$ and Matsubara frequency $\epsilon_l \equiv 2\pi l/\beta$ is written as

$$D_{ij}(\boldsymbol{q}, \epsilon_l) = \left(\delta_{ij} - \frac{q_i q_j}{q^2} \right) D(\boldsymbol{q}, \epsilon_l),$$
$$D(\boldsymbol{q}, \epsilon_l) = \{ \Pi_F(\boldsymbol{q}, \epsilon_l) + \Pi_B(\boldsymbol{q}, \epsilon_l) \}^{-1}. \quad (8)$$

Since we shall need D later in calculating $\rho(T)$, we obtain D below in the random-phase approximation as $D \simeq (\Pi^R_B + \Pi^R_F)^{-1}$. When the spin gap is sufficiently small, its effect to Π^R_F is evaluated by perturbation giving rise to a mass term as discussed above;

$$\frac{\Pi^R_F(q, \epsilon_l)}{g^2} \simeq \begin{cases} \frac{q^2}{12\pi m_F} + \sqrt{\frac{n_F}{2\pi}} \frac{|\epsilon_l|}{q} + \frac{n_F^S(T)}{m_F}, & |\epsilon_l| \ll v_F q \\ \frac{n_F}{m_F}, & |\epsilon_l| \gg v_F q \end{cases}. \quad (9)$$

We used the relation, $\mu_F \simeq \pi n_F/m_F$ and $v_F \equiv k_F/m_F$. The superfluid density of spinons is calculated for small $\Delta_{\rm SG}(\boldsymbol{k})/(k_B T)$ as

$$n_F^S(T) \simeq \frac{n_F}{2\pi} \int d\phi \left| \frac{\Delta_{\rm SG}(\boldsymbol{k})}{2k_B T} \right|^2 = \frac{n_F}{2} \left| \frac{\pi(1-\delta)\hat{\lambda}(T)}{2k_B T} \right|^2, \quad (10)$$

with $k_x/k_y = \tan\phi$ and $|\boldsymbol{k}| = k_F$. Π^R_B is given by

$$\frac{\Pi^R_B(q, \epsilon_l)}{g^2} \simeq \begin{cases} \frac{f_B(|\mu_B|)}{24\pi} \frac{q^2}{m_B} + \sqrt{\frac{n_B}{2\pi}} \frac{|\epsilon_l|}{q}, & |\epsilon_l| \ll \frac{\sqrt{n_B}}{m_B} q \\ \frac{n_B}{m_B}, & |\epsilon_l| \gg \frac{\sqrt{n_B}}{m_B} q \end{cases} \quad (11)$$

where $f_B(\epsilon) \equiv \{\exp(\beta\epsilon) - 1\}^{-1}$. Eqs.(9,11) above are obtained for small $q (\ll \pi/a)$. For large q's, they should be replaced by anisotropic expressions due to $\Delta_{\rm SG}(\boldsymbol{k})$. These anisotropy can be ignored as long as the spin gap is sufficiently small.

From the linear-response theory and Ioffe-Larkin formula [12], the dc conductivity $\sigma (\equiv \sigma_{11} = \sigma_{22})$ is expressed as

$$\sigma_{ij} = \lim_{\epsilon\to 0}\lim_{q\to 0}\frac{e^2}{-i\epsilon}\tilde{\Pi}_{ij}(\boldsymbol{q},-i\epsilon),$$
$$\tilde{\Pi}_{ij}(\boldsymbol{q},\epsilon) = \left\{\tilde{\Pi}_F^{-1}(\boldsymbol{q},\epsilon)+\tilde{\Pi}_B^{-1}(\boldsymbol{q},\epsilon)\right\}_{ij}^{-1}, \quad (12)$$

where $\tilde{\Pi}$, $\tilde{\Pi}_{F,B}$ are response functions of electron, spinon, and holon, respectively. So one has $\sigma^{-1} = \sigma_B^{-1} + \sigma_F^{-1}$.

In the spin-gap state, the spinon conductivity diverges $\sigma_F \to \infty$ due to a superflow generated by RVB condensation $\langle \lambda_{xi}\rangle \neq 0$. This is an analog of the well-known fact in the BCS theory that the electron resistivity vanishes below T_c due to a superflow generated by Cooper-pair condensation. Actually, $A^F_{\rm eff}$ has the same structure as the BCS model. Thus the total resistivity $\rho = \sigma^{-1}$ in the spin-gap state is equal to the resistivity of holons, $\rho = \sigma_B^{-1}$. Effects of spinons to ρ certainly exist and show up through the dressed propagator $D(\boldsymbol{q},\epsilon_l)$ in calculating $\tilde{\Pi}_B$.

Now we calculate the response function $\tilde{\Pi}_B$. By solving the Schwinger-Dyson equation approximately following the steps in Ref. [13], we arrive at

$$\tilde{\Pi}_{B\ ij}(0,\epsilon_l) \simeq \frac{1}{\beta}\sum_n\int\frac{d^2q}{(2\pi)^2}\frac{q_iq_j}{{m_B}^2}R_B(q,\epsilon_n;\epsilon_l) \times \frac{i\epsilon_l}{i\epsilon_l - i\epsilon_l\Gamma_B(q,\epsilon_n;\epsilon_l)-\Delta\Sigma_B(q,\epsilon_n;\epsilon_l)},$$
$$R_B(q,\epsilon_n;\epsilon_l) \equiv G_B(q,\epsilon_n)G_B(q,\epsilon_n+\epsilon_l),$$
$$G_B(q,\epsilon_n) \equiv \left(i\epsilon_n - \frac{q^2}{2m_B}+\mu_B\right)^{-1}. \quad (13)$$

$\Delta\Sigma_B(q,\epsilon_n;\epsilon_l)$, representing diagrams containing self-energy of holons, $\Sigma_B(q,\epsilon_n)$, is necessary to keep gauge invariance,

$$\Delta\Sigma_B(q,\epsilon_n;\epsilon_l) \equiv R_B^{-1}(q,\epsilon_n;\epsilon_l)\{\Sigma_B(q,\epsilon_n)G_B^2(q,\epsilon_n) - \Sigma_B(q,\epsilon_n+\epsilon_l)G_B^2(q,\epsilon_n+\epsilon_l)\}, \quad (14)$$

However, in the perturbative calculation, this combination vanishes in the dc limit by the symmetry under summations. We expect this term dose not contribute to the dc resistivity, and neglect it hereafter.

$\Gamma_B(q,\epsilon_n;\epsilon_l)$, representing vertex diagrams, contributes to σ_B,

$$\Gamma_B(q,\epsilon_n;\epsilon_l) \equiv \left(\frac{g}{m_B}\right)^2\frac{1}{\beta}\sum_{n'}\int\frac{d^2q'}{(2\pi)^2} \times\left\{\frac{\boldsymbol{q}\times(\boldsymbol{q'}-\boldsymbol{q})}{|\boldsymbol{q'}-\boldsymbol{q}|}\right\}^2\frac{\boldsymbol{q}\cdot(\boldsymbol{q'}-\boldsymbol{q})}{q^2} \times D(|\boldsymbol{q'}-\boldsymbol{q}|,\epsilon_{n'}-\epsilon_n)R_B(q',\epsilon_{n'};\epsilon_l), \quad (15)$$

where $\boldsymbol{q}\times\boldsymbol{q'} \equiv q_xq'_y - q_yq'_x$. We set q of Γ_B in Π_B to a fixed vector of typical length q_B, $q_B^2 \equiv 8\pi n_B$ [14]. We also fix the length of q' of D in (15) be q_B.

$$\Gamma_B(q_B,\epsilon_n;\epsilon_l) \simeq -\frac{g^2{q_B}^2}{8\pi^2m_B}\frac{1}{\beta}\sum_{n'}\int_{-\pi}^{\pi}d\phi\ \sin^2\phi \times D\left(q_B\sqrt{2(1-\cos\phi)},\epsilon_{n'}-\epsilon_n\right) \times\int_{|\mu_B|}^{\infty}dE\ \frac{1}{i\epsilon_{n'}-E}\frac{1}{i\epsilon_{n'}+i\epsilon_l-E}. \quad (16)$$

We consider the underdoped region, $n_B \ll n_F$ ($\delta \ll 1$), and temperatures around $\beta^{-1} \sim n_B/m_B$. Assuming that $D(q,\epsilon_l)$ in Γ_B dominates in the region near the static limit $\epsilon_l = 0$, we use the upper expressions in (9,11). In the denominator of D, the dissipation term, $\sqrt{\bar{n}/(2\pi)}|\epsilon_l|/q$, $\sqrt{\bar{n}} \equiv \sqrt{n_F}+\sqrt{n_B}$, is larger than the q^2 term, $q^2/(12\pi\bar{m})$, $\bar{m}^{-1} \equiv m_F^{-1} + f_B(|\mu_B|)/(2m_B)$, as long as $\epsilon_l \neq 0$, since their ratio is small, $(q_B^2/\bar{m})/\{\sqrt{\bar{n}}/(q_B\beta)\} \sim \mathcal{O}(\sqrt{n_B/n_F})$.

So the n'-sum is dominated at $\epsilon_{n'} = \epsilon_n$. Then we get

$$\Gamma_B(q_B,\epsilon_n;\epsilon_l) \simeq -\frac{3\bar{m}}{4\pi m_B}\frac{1}{\beta}\int_{-\pi}^{\pi}d\phi\frac{\sin^2\phi}{(1-\cos\phi)+\frac{3\bar{m}n_F^S(T)}{4m_Fn_B}} \times\left[\frac{\pi}{2\epsilon_l}\{\mathrm{sgn}(\epsilon_n+\epsilon_l)-\mathrm{sgn}(\epsilon_n)\}+\mathcal{O}(\epsilon_l^0)\right]$$
$$\simeq -\frac{1}{2\epsilon_l\tau(T)}\{\mathrm{sgn}(\epsilon_n+\epsilon_l)-\mathrm{sgn}(\epsilon_n)\} \quad (17)$$

where

$$\frac{1}{\tau(T)} \equiv \frac{3\pi\bar{m}}{2m_B}\frac{1}{\beta}\left[\left\{1+\frac{3\bar{m}n_F^S(T)}{4m_Fn_B}\right\} - \sqrt{\left\{1+\frac{3\bar{m}n_F^S(T)}{4m_Fn_B}\right\}^2-1}\right]. \quad (18)$$

To calculate $\tilde{\Pi}_{B\ ij}(0,\epsilon_l)$ we insert (17) into (13) and do the q-integral and n-sum as in (15) to get

$$\tilde{\Pi}_{B\ ij}(0,\epsilon_l) \simeq \delta_{ij}\frac{n_B}{m_B}\frac{i\epsilon_l}{\tilde{C}(T)\epsilon_l + i\tau^{-1}(T)}, \quad (19)$$

where $\lim_{\epsilon_l\to 0}\tilde{C}(T)$ is finite.

After analytic continuation $\epsilon_l > 0 \to -i\epsilon$ and using (10), we finally obtain the resistivity,

$$\rho(T) = \frac{m_B}{e^2n_B}\frac{1}{\tau(T)} \propto T\left\{1-\sqrt{\frac{3\bar{m}(1-\delta)^3}{4m_F\delta}\left|\frac{\pi\hat{\lambda}(T)}{2k_BT}\right|}\right\}+\mathcal{O}(n_F^S). \quad (20)$$

For $T^* < T < T_{\rm CSS}$, $n_F^S(T) = 0$ and this result reproduces the T-linear behavior of Ref. [5]. For T near and below T^*, one expects the behavior $|\hat{\lambda}(T)| \propto (T^*-T)^d$, with a critical exponent d, and we have $\rho(T) \propto T\{1 - c(T^*-T)^d\}$. Note also that the downward deviation of $\rho(T)$ from the T-linear behavior is reduced with increasing the doping δ. In Fig.1, we plot $\rho(T)$ of (20) with

various values of d. The MFT value $d = 1/2$ is not consistent with the experiment. To achieve smooth deviation from the T-linear behavior, one needs $d > 1$. This suggests that fluctuation effect of phases of λ_{xi} is important to obtain a realistic curve of $\rho(T)$. One could produce a reliable curve of $\rho(T)$ by inserting experimental data of $\hat{\lambda}(T)$ into (20), but the available experimental data are not enough for this purpose.

The data [6] show that one may fit ρ in a form $C_0 + C_1 T$ for $T^* < T$. This implies spinon contribution to ρ, calculated as $\sigma_F^{-1} \propto T^{4/3}/n_F$ [5] [13], is negligibly small compared with $\sigma_B^{-1} \propto T/n_B$ due to higher power in T and a small coefficient. $\sigma_F^{-1} = 0$ for $T < T^*$ as explained, but the discontinuity at $T = T^*$ in σ_F^{-1} is not observable due to its smallness.

The constant part C_0, surviving below T^*, may be attributed to scatterings of charged holons with impurities. They may be described by $H_{\rm imp} = \sum V_x b_x^\dagger b_x$, where V_x is a random potential. Actually, standard calculations show that it generates T-independent contribution to ρ, $\Delta\sigma_B^{-1} \propto 1/n_B$ at intermediate T's [15].

As T goes very low, Π_F^R is dominated by the mass term, $\Pi_F^R(q, \epsilon_l)/g^2 \simeq n_F^S(T)/m_F$, while Π_B^R does not change. The Landau damping term from holons is smaller than the mass term above. This case has been studied in Ref. [13], giving the result $\sigma_B^{-1} \propto T^2$. Thus, in $\rho(T)$, weak-localization effect by impurities, $\rho_{\rm WL} \sim C_{\rm WL}\log(T)$, dominates over inelastic scatterings by gauge bosons. This situation is in contrast with the effective gauge theory of two-dimensional electrons at half-filled Landau level, in which the transverse mode of Chern-Simons gauge field remains massless down to $T = 0$ and renormalizes $C_{\rm WL}$ [16].

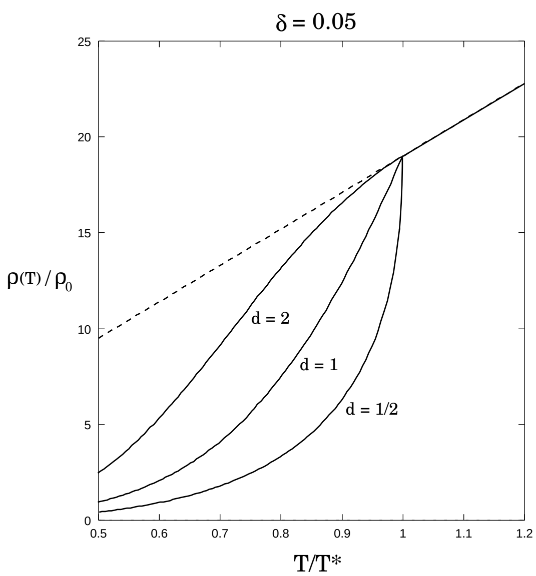

FIG. 1. Plot of the resistivity $\rho(T)$ divided by $\rho_0 \equiv 2\pi k_B T^* m_F/(e^2 n_F)$ for $\hat{\lambda}(T) \simeq \lambda_0(1 - T/T^*)^d$ ($d = 1/2, 1, 2$). For definiteness we chose $\delta = 0.05$, $\lambda_0 = 2k_B T^*/\pi$ and $3\bar{m}/(4m_F) = 0.5$.

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